

Lepton Number Violation in TeV Scale See-Saw Extensions of the Standard Model

A Ibarra¹, E Molinaro^{2,5} and S T Petcov^{3,4,6}

¹ Physik-Department T30d, Technische Universität München, James-Frank-Straße, 85748 Garching, Germany

² Centro de Física Teórica de Partículas (CFTP), Instituto Superior Técnico, Avenida Rovisco Pais 1, 1049-001, Lisboa, Portugal

³ SISSA and INFN-Sezione di Trieste, Via Bonomea 265, 34136 Trieste, Italy

⁴ Institute for the Physics and Mathematics of the Universe, University of Tokyo, Kashiwa, 277-8583, Japan

Abstract.

The low-energy neutrino physics constraints on the TeV scale type I see-saw scenarios of neutrino mass generation are revisited. It is shown that lepton charge (L) violation, associated to the production and decays of heavy Majorana neutrinos N_j having masses in the range of $M_j \sim (100 \div 1000)$ GeV and present in such scenarios, is hardly to be observed at ongoing and future particle accelerator experiments, LHC included, because of very strong constraints on the parameters and couplings responsible for the corresponding $|\Delta L| = 2$ processes. If the heavy Majorana neutrinos N_j are observed and they are associated only with the type I mechanism, they will behave effectively like pseudo-Dirac fermions. Conversely, the observation of effects proving the Majorana nature of N_j would imply that these heavy neutrinos have additional relatively strong couplings to the Standard Model particles or that light neutrino masses compatible with the observations are generated by a mechanism other than see-saw (e.g., radiatively at one or two loop level) in which the heavy Majorana neutrinos N_j are nevertheless involved.

1. Introduction

Neutrino oscillation experiments [1] have revealed undeniable evidence for flavour non-conservation in the lepton sector. Flavour neutrino eigenstates, $\nu_{\ell L}$ ($\ell = e, \mu, \tau$), oscillate, changing their lepton flavour quantum number [2, 3], due to non-zero neutrino masses and neutrino mixing. The latter is parametrized by the well known Pontecorvo-Maki-Nakagawa-Sakata (PMNS) neutrino mixing matrix [4].

All present neutrino oscillation data can be described assuming 3-flavour neutrino mixing in vacuum. The number of massive neutrinos can, in general, be bigger than 3, if, for instance, there exist right-handed (RH) sterile neutrinos [5] and they mix with the left-handed (LH) flavour neutrinos. It follows from the existing data that at least 3 of the neutrinos must be light, with masses $m_1 \neq m_2 \neq m_3$, $m_{1,2,3} \lesssim 1$ eV and at least two of the m_j different from zero. At present there are no compelling experimental evidences for the existence of more than 3 light neutrinos.

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⁶ Also at: Institute of Nuclear Research and Nuclear Energy, Bulgarian Academy of Sciences 1784 Sofia, Bulgaria.

The smallness of neutrino masses, $m_j/m_e \lesssim 10^{-6}$, where m_e is the electron mass, suggests the presence of a new fundamental scale in particle physics and thus to new physics beyond the Standard Model (SM). Among the possible extensions of the SM which can explain neutrino mass generation, a well known proposal is the see-saw mechanism [6]. In its simplest formulation, the type I see-saw scenario, at least two “heavy” $SU(2)_L \times U(1)_Y$ singlet RH neutrinos ν_{aR} are introduced in the theory. Within the see-saw framework, the latter are assumed to possess a Majorana mass term as well as Yukawa type couplings with the Standard Model lepton and Higgs doublets.

Below the electroweak (EW) symmetry breaking scale, the light neutrino masses arise from the Lagrangian:

$$\mathcal{L}_\nu = -\overline{\nu_{\ell L}}(M_D)_{\ell a}\nu_{aR} - \frac{1}{2}\overline{\nu_{aL}^C}(M_N)_{ab}\nu_{bR} + \text{h.c.}, \quad (1)$$

where $\nu_{aL}^C \equiv C\overline{\nu_{aR}}^T$ ($a = 1, \dots, k$), C being the charge conjugation matrix, $M_N = (M_N)^T$ is the $k \times k$ Majorana mass matrix of the RH neutrinos, and M_D is the $3 \times k$ neutrino Dirac mass matrix, which is generated by the matrix of neutrino Yukawa couplings after the EW symmetry breaking: $M^D = v\lambda$, $\lambda_{\ell a}$ being the matrix of neutrino Yukawa couplings and $v = 174$ GeV being the v.e.v. of the SM Higgs doublet. The matrices M_N and M_D are complex, in general. A Majorana mass term m_ν for the active left-handed neutrinos is indeed generated by the interplay between the Dirac mass term and the Majorana mass term of the heavy Majorana neutrinos. The well known see-saw mass relation is in this case: $m_\nu \cong -M_D M_N^{-1} (M_D)^T$.

In grand unified theories, M_D is typically of the order of the charged fermion masses. In $SO(10)$ theories, for instance, M_D coincides with the up-quark mass matrix. Taking indicatively $m_\nu \sim 0.05$ eV, $M_D \sim 100$ GeV, one finds $M_N \sim 2 \times 10^{14}$ GeV, which is close to the scale of unification of the electroweak and strong interactions, $M_{GUT} \cong 2 \times 10^{16}$ GeV. In GUT theories with RH neutrinos one finds that indeed the heavy Majorana neutrinos naturally obtain masses which are by few to several orders of magnitude smaller than M_{GUT} . The estimate of M_N thus obtained is effectively based on the assumption that the neutrino Yukawa couplings are large: $|\lambda_{\ell a}| \sim 1$. The alternative possibility is to have heavy Majorana neutrino masses in the range of $\sim (100 \div 1000)$ GeV, i.e. TeV scale see-saw generation of neutrino masses. This possibility has been recently revisited by several authors (see e.g. [7] and references quoted therein).

Following the discussion reported in [7], the low-energy neutrino physics constraints on the TeV scale type I see-saw scenarios are reviewed. In particular, it is analyzed in detail, within three different frameworks relying on the type I see-saw extension of the SM, the possibility of observing $|\Delta L| = 2$ processes at current and future particle accelerators, LHC included, due to the production and decays of the heavy Majorana neutrinos.

2. Non-unitarity effects, see-saw mechanism and $(\beta\beta)_{0\nu}$ -decay

Following the formalism developed in [7], after the diagonalization of the full neutrino mass matrix in (1), the charged current (CC) and neutral current (NC) weak interactions involving the light Majorana neutrinos χ_j with definite mass m_j are

$$\mathcal{L}_{CC}^\nu = -\frac{g}{\sqrt{2}}\bar{\ell}\gamma_\alpha\nu_{\ell L}W^\alpha + \text{h.c.} = -\frac{g}{\sqrt{2}}\bar{\ell}\gamma_\alpha((1+\eta)U)_{\ell i}\chi_{iL}W^\alpha + \text{h.c.}, \quad (2)$$

$$\mathcal{L}_{NC}^\nu = -\frac{g}{2c_w}\overline{\nu_{\ell L}}\gamma_\alpha\nu_{\ell L}Z^\alpha = -\frac{g}{2c_w}\overline{\chi_{iL}}\gamma_\alpha\left(U^\dagger(1+\eta+\eta^\dagger)U\right)_{ij}\chi_{jL}Z^\alpha. \quad (3)$$

The unitary matrix U diagonalizes the light Majorana mass matrix, $m_\nu = U^*\text{diag}(m_1, m_2, m_3)U^\dagger$. In the basis in which the three charged lepton fields are equal to the respective mass eigenstates, U corresponds to the PMNS neutrino mixing matrix. The hermitian matrix η , therefore, parametrizes the deviation from unitarity of the neutrino mixing matrix due to the mixing of

the flavour neutrino fields $\nu_{\ell L}$ with the heavy RH fields in the neutrino mass Lagrangian (1). As shown below, η depends on a combination of the see-saw parameters responsible for lepton flavour violating processes. Notice that in deriving the expressions of the CC and NC weak interactions it is implicitly assumed that $\eta_{ij} \ll 1$. Indeed, for SM singlet fermions with masses above the EW symmetry breaking scale, i.e. bigger than ~ 100 GeV, the combined data on neutrino oscillation experiments and EW interaction processes provide the upper bound [8, 9]: $|\eta_{ij}| < 6 \times 10^{-3}$.

The charged current and the neutral current interactions of the heavy Majorana neutrino mass eigenstates N_j with W^\pm and Z^0 formally read [7]:

$$\mathcal{L}_{CC}^N = -\frac{g}{2\sqrt{2}} \bar{\ell} \gamma_\alpha (RV)_{\ell k} (1 - \gamma_5) N_k W^\alpha + \text{h.c.}, \quad (4)$$

$$\mathcal{L}_{NC}^N = -\frac{g}{2c_w} \bar{\nu}_{\ell L} \gamma_\alpha (RV)_{\ell k} N_k Z^\alpha + \text{h.c.}, \quad (5)$$

where the unitary matrix V diagonalizes the RH neutrino mass term in (1), $M_N \simeq V^* \text{diag}(M_1, \dots, M_k) V^\dagger$, and $R \equiv (M_D M_N^{-1})^*$. The experimental limits on the deviations from unitarity in the neutrino mixing matrix affect the couplings $(RV)_{\ell j}$. Indeed, in this class of models [7]: $\eta = -\frac{1}{2} R R^\dagger$.

It is important to remark that, apart from non-unitarity effects, further constraints on $(RV)_{\ell j}$ and/or a particular combination of them can be obtained at low energy considering *i*) the scale of the symmetric matrix $m_\nu = -M_D M_N^{-1} (M_D)^T \simeq -R^* M_N R^\dagger \lesssim 1$ eV, resulting from the see-saw mechanism of generation of neutrino masses, and *ii*) from the experimental upper bound on the effective Majorana mass, $|(m_\nu)_{ee}|$ [10], which enters in the expression of the $(\beta\beta)_{0\nu}$ -decay rate of even-even nuclei (see e.g. [11, 12]). Indeed, for the heavy Majorana neutrinos N_k with masses $M_k \sim M_R \geq 100$ GeV, barring “accidental” cancellations or extreme fine-tuning (at the level of $\sim 10^9$), the constraint *i*) implies:

$$|(RV)_{\ell j}| \lesssim 3 \times 10^{-6} \left(\frac{100 \text{ GeV}}{M_R} \right)^{1/2}. \quad (6)$$

The estimate of $|(RV)_{\ell j}|$ thus obtained makes the heavy Majorana neutrinos N_j practically unobservable at particle accelerators (see e.g. [13]). In order for the CC and NC couplings of the heavy Majorana neutrinos N_j to W^\pm and Z^0 , eqs (4) and (5), to be sufficiently large, that is $|(RV)_{\ell j}| \approx 10^{-2}$, so that the see-saw mechanism could be partially or completely tested in experiments at the currently operating and planned future accelerators (LHC included), the small size of the m_ν matrix elements should be due to strong mutual compensation between the parameters that enter in the see-saw mass relation. In general, such cancellations arise naturally from symmetries in the lepton sector, corresponding, for instance, to the conservation of some additive lepton charge \hat{L} (see, e.g. [11, 14, 15, 16]). In the limit in which \hat{L} is exact, the RH neutrinos ν_{aR} ($a = 1, \dots, k$) should be Dirac particles, which is possible for all heavy neutrinos only if the number of the RH singlet neutrino fields is even: $k = 2q$, $q = 1, 2, \dots$. If their number is odd, barring again “accidental” cancellations, some (odd number) of the discussed heavy Majorana neutrinos will have strongly suppressed couplings to the W^\pm and charged leptons and will be practically unobservable in the current and the future planned accelerator experiments.

The resulting mass spectrum will depend on the assumed total lepton charge \hat{L} . Indeed, if L_i ($i = e, \mu, \tau, \dots, k$) are the ordinary lepton charges of the fermion fields, in the basis in which the charged lepton mass matrix is diagonal, the conserved lepton charge \hat{L} takes the general form:

$$\hat{L} = \sum_{i=e,\mu,\tau,\dots,k} (-1)^{n_i} a_i L_i, \quad (7)$$

where $n_i = 0, 1$ and $a_i = 0, 1$ for $i = e, \mu, \tau, \dots, k$, with at least one $a_i \neq 0$. In this case, it can be shown [14] that the number of massive Dirac neutrinos and massless neutrinos in the theory is, respectively, $\min(n_+(\hat{L}), n_-(\hat{L}))$ and $|n_+(\hat{L}) - n_-(\hat{L})|$, where $n_+(\hat{L})$ ($n_-(\hat{L})$) is the number of ordinary lepton charges in (7) which enter with the plus (minus) sign.

The correct light Majorana neutrino mass spectrum can be generated by small perturbations that violate the corresponding symmetry \hat{L} . In this case, each massive Dirac neutrino is split into two Majorana neutrinos with close but different masses, i.e. the Dirac neutrinos become pseudo-Dirac fermions [17, 18]. Typically, for $|(RV)_{\ell j}| \sim 10^{-3}$ (10^{-4}), the expected splitting between the masses of the two heavy Majorana neutrinos forming a pseudo-Dirac pair, is roughly 1 (100) MeV for masses of the order of 100 (1000) GeV. Such small mass difference will not be observable at LHC and planned future accelerator experiments, thus preventing any possible experimental test at collider of the Majorana nature of the heavy Majorana neutrinos of the type I see-saw mechanism.

The other low energy constraint to the see-saw parameter space, strictly related to the Majorana nature of the heavy and light Majorana neutrinos, is provided by $(\beta\beta)_{0\nu}$ -decay experiments [10]. In this case, in addition to the standard contribution due to the light Majorana neutrino exchange, the $(\beta\beta)_{0\nu}$ -decay effective Majorana mass $|(m_\nu)_{ee}|$ receives a contribution from the exchange of the heavy Majorana neutrinos N_k :

$$|(m_\nu)_{ee}| \cong \left| \sum_i (U_{\text{PMNS}})_{ei}^2 m_i - \sum_k F(A, M_k) (RV)_{ek}^2 M_k \right|, \quad (8)$$

where $F(A, M_k)$ is a known real (positive) function of the atomic number A of the decaying nucleus and of the mass M_k of N_k [19, 20, 21]. For the range $M_k \sim (100 \div 1000)$ GeV of interest, one has with good approximation (see, e.g. [20, 21, 22, 23]) $F(A, M_k) \simeq (M_a/M_k)^2 f(A)$, with $M_a \simeq 0.9$ GeV and $f(A) \simeq \mathcal{O}(10^{-2})$.

3. The case of a broken lepton charge \hat{L}

It is shown in the following that, within the type I see-saw scenario of generation of light neutrino masses, if $m_\nu \neq 0$ arises as result of the breaking of a global symmetry corresponding to some conserved lepton charge \hat{L} , the Majorana nature of the heavy Majorana neutrinos could be hardly tested at collider physics. As stated before, the reason is that if the heavy Majorana fields involved in the see-saw mechanism of LH neutrino mass generation, have couplings to charged leptons and weak gauge bosons sufficiently large to produce and observe them at collider, they would behave effectively as a pseudo-Dirac fermion and all lepton number violating processes, e.g. same sign di-muon production in proton-proton collisions, would be strongly suppressed.

Consider for simplicity the case of three RH neutrinos ν_{aR} ($a = 1, 2, 3$) and a conserved lepton charged $\hat{L} = L_e + L_\mu + L_\tau + L_3 - L_1$ (see e.g. [7]). A possible choice of the Dirac and Majorana mass matrix which preserves \hat{L} is:

$$M_D = \begin{pmatrix} 0 & 0 & m_{e3}^D \\ 0 & 0 & m_{\mu 3}^D \\ 0 & 0 & m_{\tau 3}^D \end{pmatrix}, \quad M_N = \begin{pmatrix} M_{11} & 0 & 0 \\ 0 & 0 & M_{23} \\ 0 & M_{23} & 0 \end{pmatrix}, \quad (9)$$

where, without loss of generality, all the parameters in M_D and M_N are taken real and positive. As a consequence of the simplifying choice made, $m_{\ell 1}^D = 0$, $\ell = e, \mu, \tau$, the Majorana field $N_1 = (\nu_{1R} + \nu_{1L}^C)/\sqrt{2}$ is decoupled and does not take part in the see-saw mechanism of generation of LH neutrino masses. In this case $n_+(\hat{L}) = 4$ and $n_-(\hat{L}) = 1$ and the spectrum of the theory contains 3 massless neutrinos and 1 massive Dirac fermion. The latter is given by the

combination $N_D = (N_2 + N_3)/\sqrt{2}$ and has a mass $M = M_{23} > 0$, where the two heavy Majorana neutrinos N_2 and N_3 have the same mass $M_2 = M_3 = M_{23}$ and satisfy the Majorana conditions $C\overline{N}_k^T = \rho_k N_k$, $k = 2, 3$, with $\rho_2 = -1$ and $\rho_3 = +1$.

In order to have a light neutrino mass spectrum compatible with the observations, it is necessary to softly break \hat{L} . Taking $(M_D)_{\ell 2} \equiv m_{\ell 2}^D \neq 0$, it is easy to show [7] that two of the light neutrinos acquire a finite mass, $0 < m_2 < m_3$, while the Dirac fermion N_D is split into two different Majorana fields N_2 and N_3 with $\Delta M \equiv M_3 - M_2 = m_3 - m_2 = 2A/M_{23}$, where A depends explicitly on the lepton charge breaking parameters $m_{\ell 2}^D$: $A = m_{e 2}^D m_{e 3}^D + m_{\mu 2}^D m_{\mu 3}^D + m_{\tau 2}^D m_{\tau 3}^D$. The light neutrino mass scale, $m_3 - m_2 \lesssim 1$ eV, imposes strong constraints on the symmetry breaking parameters. More explicitly, for $M_{23} \approx 100$ GeV one has: $|m_{\ell 2}^D m_{\ell 3}^D| \lesssim 10^{-7} \text{ GeV}^2$ ($\ell, \ell' = e, \mu, \tau$). Because of the small mass splitting ΔM , the heavy Majorana neutrinos $N_{2,3}$ effectively behave in the collider physics as a pseudo-Dirac fermion $N_{PD} = (N_2 + N_3)/\sqrt{2}$.

Consider now the process of same sign di-muon production in proton-proton collisions, assuming that one of the muons, say μ^- , is produced together with real or virtual $N_{2,3}$ in the decay of a virtual W^- , while the second μ^- originates from the decay $N_{2,3} \rightarrow W^+ \mu^-$, with virtual or real W^+ which decays further into, e.g. two hadronic jets. In order to collect a sufficiently large number of events, it is necessary to have sizable couplings of the heavy Majorana neutrinos to the muon field. The latter can be easily computed [7]: $(RV)_{\mu 2} = m_{\mu 3}^D \cos \theta / M_{23}$ and $(RV)_{\mu 3} = m_{\mu 3}^D \sin \theta / M_{23}$, with $\tan^2 \theta = M_2 / M_3$. In the present scenario $\theta \simeq \pi/4$ because N_2 and N_3 form a pseudo-Dirac fermion. If, for instance, $M_{23} \approx 100$ GeV and $|(RV)|_{\mu 2,3} \approx 10^{-2}$, one obtains $m_{\mu 3}^D \approx 1$ GeV and \hat{L} charge breaking parameter $m_{\mu 2}^D \approx 100$ eV, due to the constraint which arises from the small neutrino mass scale (see the discussion above). The relevant part for the amplitude of the process under discussion is [7]:

$$A(p p \rightarrow \mu^- \mu^- 2\text{jets } X) \propto \frac{(m_{\mu 3}^D)^2}{M_{23}^2} \frac{M_2 M_3}{M_3 + M_2} \frac{M_3^2 - M_2^2 - i(\Gamma_3 M_3 - \Gamma_2 M_2)}{(p^2 - M_3^2 + i\Gamma_3 M_3)(p^2 - M_2^2 + i\Gamma_2 M_2)}. \quad (10)$$

Taking into account that $M_2 M_3 \simeq M_{23}^2$ and $\Gamma_{2(3)} \propto \sum_{\ell} |(RV)_{\ell 2,3}|^2 G_F M_{2(3)}^3$, it is easy to show that the amplitude for this lepton number violating process is proportional to the factor $\Delta M / M_{23}$. As discussed previously, in the pseudo-Dirac scenario of interest, which is consequence of an almost conserved lepton charge present in the theory, the heavy Majorana neutrino mass splitting is strongly bounded by the light neutrino mass scale, $\Delta M \lesssim 1$ eV, thus making the same sign di-muon production at LHC unobservable.

4. \hat{L} is not conserved but $m_\nu = 0$ at leading order

A simple realization of this scenario is obtained from the mass textures reported in eq. (9), in which the Majorana mass matrix, M_N , is conveniently modified with the \hat{L} -breaking mass term: $(M_N)_{33} = M_{33} > 0$. In this case, even if lepton number is not conserved, at leading order one has $m_\nu = M_D M_N^{-1} M_D^T = 0$ and, therefore, the tree-level constraints on the see-saw parameter space from the light neutrino mass scale do not apply here: light neutrino masses, compatible with neutrino data, can be generated at higher order (one or two loops). This in turn could lead to sufficiently large $N_{2,3}$ production rates at LHC to make the observation of the two heavy Majorana neutrinos possible. Indeed, in such scenario $\Delta M = M_3 - M_2 = M_{33}$ and, in principle, one can observe the same sign di-muon production at LHC for a sufficiently large mass splitting (see eq. (10)). Anyway, it should be clear from the preceding discussion that all $|\Delta L| = 2$ Majorana type effects should vanish in the limit of $M_{33} = 0$.

As discussed before, in such class of models, the heavy Majorana neutrinos might provide a sizable contribution to the $(\beta\beta)_{0\nu}$ -decay rate. Indeed, from eq. (8), neglecting the standard

contribution due to the tree-level exchange of light Majorana neutrino mass eigenstates, the effective Majorana mass $|(m_\nu)_{ee}|$ results proportional to the mass splitting ΔM [7]:

$$|(m_\nu)_{ee}| \cong \left| \frac{(m_{e3}^D)^2}{M_{23}^2} f(A) M_a^2 \frac{M_3 - M_2}{M_2 M_3} \right| \cong \left| \frac{(m_{e3}^D)^2}{M_{23}^2} f(A) \frac{M_a^2}{M_{23}^2} M_{33} \right|. \quad (11)$$

If $(m_{e3}^D)^2/M_{23}^2$ will be determined from an independent measurement, the bound on $|(m_\nu)_{ee}|$ will lead to a strong constraint on M_{33}/M_{23} . Taking e.g. $|(m_\nu)_{ee}| \lesssim 1$ eV, $f(A) = 0.078$ (corresponding to ^{76}Ge) and the maximal value of $(m_{e3}^D)^2/M_{23}^2$ allowed by the non-unitarity constraints (see [8]), 8×10^{-3} , one finds: $M_{33} \lesssim 1.8 \times 10^{-5} M_{23}(M_{23}/M_a)$. For $M_{23} = 100$ GeV this implies $M_{33} \lesssim 2 \times 10^{-3} M_{23} \cong 0.2$ GeV $\ll M_{23}$. Such a small mass difference would render the Majorana-type effects associated with $N_{2,3}$ hardly observable. If, however, $|(m_{e3}^D/M_{23})^2| \lesssim 1.6 \times 10^{-6}$, $M_{33} \lesssim M_{23}$ and a signature of the process $pp \rightarrow \mu^- \mu^- 2 \text{ jets } X$, generated by the production and decay of real or virtual $N_{2,3}$, is detectable at LHC.

5. The extreme fine-tuning case

Consider now the type I see-saw Lagrangian in which the RH neutrino masses are given between $(100 \div 1000)$ GeV and the see-saw mechanism provides the observed LH neutrino mass scale and mixing without assuming any softly broken conserved lepton charge in the theory. This possibility requires, in general, a huge tuning of parameters. Namely, demanding $(M_D)_{ij} \sim \mathcal{O}(1 \text{ GeV})$ and $M_j \sim \mathcal{O}(100 \text{ GeV})$ requires a tuning of one part in 10^9 in order to produce a neutrino mass $m_i \sim \mathcal{O}(10^{-2} \text{ eV})$ [7].

Unlike the case discussed earlier, now the see-saw parameter space is constrained by both the small light neutrino mass scale, $m_\nu \lesssim 1$ eV, and the upper bound on the effective Majorana mass in $(\beta\beta)_{0\nu}$ -decay experiments. Consider, for instance, the two heavy Majorana neutrinos mass eigenstates N_2 and N_3 with masses $M_2 > 0$ and $M_3 \equiv M_2(1 + z)$, $z > 0$, and satisfying standard Majorana conditions, $C\overline{N}_{2,3}^T = N_{2,3}$. For large neutrino Yukawa couplings, the CC and NC weak interactions of N_2 and N_3 with charged leptons and gauge bosons, eqs (4) and (5), satisfy, with a good approximation, the relation [7]: $(RV)_{\ell 3} \simeq i(RV)_{\ell 2}/\sqrt{1+z}$ ($\ell = e, \mu, \tau$), with $|(RV)_{\mu 2}| \approx |(RV)_{e 2}|$. From eq. (8), neglecting again the contribution due to the exchange of the light Majorana neutrino mass eigenstates, the effective Majorana mass is given by

$$|(m_\nu)_{ee}| \cong 2z |(RV)_{e 2}|^2 \frac{M_a^2}{M_2} f(A). \quad (12)$$

Assuming production and detection of the RH Majorana fields at collider, i.e. $|(RV)_{\ell 2}| \approx 10^{-2}$ ($\ell = e, \mu$), one has that $|(m_\nu)_{ee}| \lesssim 0.2$ eV for $z \lesssim 10^{-3}$ (10^{-2}) and $M_1 \approx 100$ (1000) GeV. Therefore, the non-observation of the $(\beta\beta)_{0\nu}$ -decay requires in this scenario a degeneracy in the right-handed neutrino masses of at least one per cent. As discussed above, the cross section for same sign di-muon production in proton-proton collisions is proportional to the mass difference of the right-handed neutrinos. Thus, even in this extremely fined-tuned scenario, the Majorana nature of the right-handed neutrinos will be difficult to probe at colliders.

6. Conclusions

The see-saw mechanism provides a natural explanation for the smallness of neutrino masses. In such scenario, a new mass scale is introduced in the theory, which is linked to the light neutrino mass scale and can in principle be accessible to present and forthcoming particle physics accelerators, LHC included. In its simplest formulation, the type I see-saw scenario, the Standard Model (SM) is extended with at least two “heavy” Majorana neutrinos, which are singlets of $SU(2)_L \times U(1)_Y$ and are coupled to the SM lepton and Higgs doublets. It is

shown that, if the type I see-saw scenario provides the correct mechanism of generation of light neutrino masses and the Majorana fields N_j have masses in the range $M_j \sim (100 \div 1000)$ GeV, the physical effects associated with the Majorana nature of these heavy neutrinos N_j , are so small that they are unlikely to be observable in the currently operating and future planned accelerator experiments (including LHC). This is a consequence of the existence of very strong constraints on the parameters and couplings, responsible for the corresponding $|\Delta L| = 2$ processes in which N_j are involved, and/or on the couplings of N_j to the weak W^\pm and Z^0 bosons. The strongest constraints are provided by the experimental upper limit on *i*) the light neutrino mass matrix, $|(m_\nu)_{ll'}| \lesssim 1$ eV, $l, l' = e, \mu, \tau$ and *ii*) the effective Majorana mass $|(m_\nu)_{ee}|$ obtained in the $(\beta\beta)_{0\nu}$ -decay experiments. The latter, is sensitive to the heavy Majorana neutrino mass scale and splitting(s).

As consequence, it is shown that charged and neutral current interactions of the heavy Majorana fields N_j with the Standard Model charged leptons and neutrinos are extremely suppressed, unless the heavy Majorana neutrinos form a pseudo-Dirac pair. Therefore, the Majorana nature of the heavy neutrinos will not be detected in collider experiments: either the production cross section is highly suppressed or the heavy neutrinos behave to a high level of precision as Dirac fermions.

The suppression of the $|\Delta L| = 2$ processes can be avoided only if there exist additional TeV scale interaction terms in the Lagrangian between the heavy Majorana neutrinos and the Standard Model particles. If this is the case, the production cross section of heavy neutrinos will not necessarily be suppressed, while their charged and neutral current interactions with the Standard Model charged leptons and neutrinos can still be tiny.

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References

- [1] B.T. Cleveland *et al.*, *Astrophys. J.* **496**, 505 (1998); Y. Fukuda *et al.* [Kamiokande Collaboration], *Phys. Rev. Lett.* **77**, 1683 (1996); J.N. Abdurashitov *et al.*, *Phys. Rev. C* **80**, 015807 (2009); P. Anselmann *et al.*, *Phys. Lett. B* **285** (1992) 376; W. Hampel *et al.*, *Phys. Lett. B* **447** (1999) 127; M. Altmann *et al.*, *Phys. Lett. B* **616** (2005) 174; S. Fukuda *et al.* [Super-Kamiokande Collaboration], *Phys. Lett. B* **539** (2002) 179; Q.R. Ahmad *et al.* [SNO Collaboration], *Phys. Rev. Lett.* **87** (2001) 071301 and **89** (2002) 011301; Y. Fukuda *et al.* [Super-Kamiokande Collaboration], *Phys. Rev. Lett.* **81** (1998) 1562; Y. Ashie *et al.* [Super-Kamiokande Collaboration], *Phys. Rev. Lett.* **93** (2004) 101801; K. Eguchi *et al.* [KamLAND Collaboration], *Phys. Rev. Lett.* **90** (2003) 021802; T. Araki *et al.*, *Phys. Rev. Lett.* **94** (2005) 081801; C. Arpesella *et al.*, *Phys. Lett. B* **658** (2008) 101; *Phys. Rev. Lett.* **101** (2008) 091302; M. H. Ahn *et al.* [K2K Collaboration], *Phys. Rev. D* **74** (2006) 072003; D.G. Michael *et al.* [MINOS Collaboration], *Phys. Rev. Lett.* **97** (2006) 191801; P. Adamson *et al.*, *Phys. Rev. Lett.* **101** (2008) 131802.
- [2] B. Pontecorvo, *Zh. Eksp. Teor. Fiz. (JETP)* **33** (1957) 549 and **34** (1958) 247.
- [3] Z. Maki, M. Nakagawa and S. Sakata, *Prog. Theor. Phys.* **28** (1962) 870.
- [4] C. Amsler *et al.* [Particle Data Group], *Phys. Lett. B* **667** (2008) 1.
- [5] B. Pontecorvo, *Sov. Phys. JETP* **26** (1968) 984. [*Zh. Eksp. Teor. Fiz.* **53**, 1717 (1967)].
- [6] P. Minkowski, *Phys. Lett. B* **67** (1977) 421; M. Gell-Mann, P. Ramond and R. Slansky, *Proceedings of the Supergravity Stony Brook Workshop*, New York 1979, eds. P. Van Nieuwenhuizen and D. Freedman; T. Yanagida, *Proceedings of the Workshop on Unified Theories and Baryon Number in the Universe*, Tsukuba, Japan 1979, eds. A. Sawada and A. Sugamoto; R. N. Mohapatra and G. Senjanovic, *Phys. Rev. Lett.* **44** (1980) 912.
- [7] A. Ibarra, E. Molinaro and S. T. Petcov, *JHEP* **1009** (2010) 108 [arXiv:1007.2378 [hep-ph]].

- [8] S. Antusch, J. P. Baumann and E. Fernandez-Martinez, Nucl. Phys. B **810** (2009) 369.
- [9] S. Antusch *et al.*, JHEP **0610** (2006) 084.
- [10] See, e.g. C. Aalseth *et al.*, arXiv:hep-ph/0412300; C.E. Aalseth *et al.* [IGEX Collaboration], Phys. Atomic Nuclei **63** (2000) 1225; H. V. Klapdor-Kleingrothaus *et al.*, Phys. Lett. B **586** (2004) 198; A. S. Barabash [NEMO Collaboration], arXiv:0807.2336 [nucl-ex]; C. Arnaboldi *et al.* [CUORICINO Collaboration], Phys. Rev. C **78** (2008) 035502.
- [11] S. M. Bilenky and S. T. Petcov, Rev. Mod. Phys. **59** (1987) 671.
- [12] S.M. Bilenky, S. Pascoli and S.T. Petcov, Phys. Rev. D **64** (2001) 053010; S.T. Petcov, Phys. Scripta T **121** (2005) 94; S. Pascoli and S. T. Petcov, Phys. Rev. D **77** (2008) 113003.
- [13] A. Atre *et al.*, JHEP **0905** (2009) 030.
- [14] C.N. Leung and S.T. Petcov, Phys. Lett. B **125** (1983) 461.
- [15] D. Wyler and L. Wolfenstein, Nucl. Phys. B **218** (1983) 205.
- [16] G. C. Branco, W. Grimus and L. Lavoura, Nucl. Phys. B **312** (1989) 492.
- [17] L. Wolfenstein, Nucl. Phys. B **186** (1981) 147.
- [18] S.T. Petcov, Phys. Lett. B **110** (1982) 245.
- [19] A. Halprin, S.T. Petcov and S.P. Rosen, Phys. Lett. B **125** (1983) 335.
- [20] J. Vergados, Nucl. Phys. B **218** (1983) 109.
- [21] W.C. Haxton and J. Stephenson, Prog. Part. Nucl. Phys. **12** (1984) 409.
- [22] H. Paes *et al.*, Phys. Lett. B **498** (2001) 35.
- [23] M. Blennow *et al.*, JHEP **1007** (2010) 096 [arXiv:1005.3240 [hep-ph]].